Deeply virtual Compton scattering in next-to-leading order.

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Abstract

We study the amplitude of deeply virtual Compton scattering in next-to-leading order of perturbation theory including the two-loop evolution effects for different sets of skewed parton distributions (SPDs). It turns out that in the minimal subtraction scheme the relative radiative corrections are of order 20-50%. We analyze the dependence of our predictions on the choice of SPD, that will allow to discriminate between possible models of SPDs from future high precision experimental data, and discuss shortly theoretical uncertainties induced by the radiative corrections.

Keywords: deeply virtual Compton scattering, skewed parton distribution, next-to-leading order corrections

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Virtual Compton scattering (VCS) in a light cone dominated region, — usually referred to as the deeply virtual Compton scattering (DVCS) [1]-[3], — is a favourable process to get access to the so-called skewed parton distributions (SPDs) [1]-[6]. The latter can be viewed as a generalization of the conventional parton densities as measured in deep inelastic scattering (DIS). They contain complementary information about the internal structure of hadrons, e.g. the total angular momentum carried by quarks and gluons [2].

The study of SPDs promises to become one of the main issues of hadronic physics for the next decade. This marks general trends in the development of experimental techniques towards exclusive reactions at the facilities like CERN and DESY. In general, it will be a delicate task to measure DVCS, since it is contaminated by the Bethe-Heitler (BH) process. However, it is rather encouraging that a first DVCS signal has already been detected by the ZEUS collaboration [7]. In addition, azimuthal, spin and charge asymmetries [8] allow to get access to the interference term between DVCS and BH processes and this should allow to separate real and imaginary parts of the DVCS amplitude. Such measurements would pin down plausible shapes of SPDs.

The theoretical goal is, therefore, to push our understanding of SPDs to the level reached nowadays for the usual parton distribution functions. One of the most important questions is the Q^2-evolution of SPDs. In earlier papers we gave a complete solution at the NLO level [9]. In the present paper we study the rôle of the NLO radiative corrections for the DVCS amplitude quantitatively in the region $x_{\rm Bj} > 10^{-2}$ including the effects of scaling violations. Our aim is to decide whether an extension to NNLO is necessary to obtain a stable interpretation of a given set of DVCS data. For a first rough comparison with experimental measurements NLO is probably fine but to judge the potential of SPDs measurement the precise size of the corrections as well as the uncertainties involved have to be known.

The VCS amplitude is defined in terms of the time ordered product of two electromagnetic currents

$$T_{\mu\nu}(P, q_1, q_2) = i \int dx e^{ixq} \langle P_2, S_2 | T \left\{ j_\mu(x/2) j_\nu(-x/2) \right\} | P_1, S_1 \rangle. \tag{1}$$

Here P_1 and P_2 are the momenta of the initial and the final hadrons, respectively. The incoming photon with momentum q_1 has a large virtuality. The scaling variables, ξ and η, which allow to describe different "two-photon" processes in the light-cone dominated region are introduced as follows [1]

$$\xi \equiv \frac{Q^2}{Pq}, \quad \eta \equiv \frac{\Delta q}{Pq}, \quad \text{where} \tag{2}$$
$$Q^2 = -q^2 = -\frac{1}{4}(q_1 + q_2)^2, \quad P = P_1 + P_2, \quad \Delta = P_1 - P_2 = q_2 - q_1.$$

In DVCS kinematics $-q_1^2 \gg m^2_{\rm hadron}$, $q_2^2 = 0$ and we have in fact only one scaling variable ξ since $\eta = \xi \left(1 - \frac{\Delta^2}{4Q^2}\right) \approx \xi$. From the experimental point of view it is more appropriate to work with

the variables $-q_1^2$ and $x_{\rm Bj} \equiv -q_1^2/2(P_1 q_1)$, which are related to the variables (2) by

$$Q^2 = -\frac{1}{2} q_1^2 \left(1 - \frac{\Delta^2}{2 q_1^2}\right) \approx -\frac{1}{2} q_1^2, \quad \xi = \frac{x_{\rm Bj}\left(1 - \frac{\Delta^2}{2q_1^2}\right)}{2 - x_{\rm Bj}\left(1 + \frac{\Delta^2}{q_1^2}\right)} \approx \frac{x_{\rm Bj}}{2 - x_{\rm Bj}}. \tag{3}$$

In the kinematical region we are interested in there are two leading twist contributions

$$T_{\mu\nu}(P, q_1, q_2) = -\tilde{g}^T_{\mu\nu} \mathcal{F}_1(\xi, \eta = \xi, Q^2, \Delta^2) + i \tilde{\epsilon}_{\mu\nu q P} \frac{1}{Pq} \mathcal{G}_1(\xi, \eta = \xi, Q^2, \Delta^2) + \cdots, \tag{4}$$

where the ellipsis stand for a leading twist-two contribution coming from longitudinally polarized photons which, however, does not appear in the DVCS cross section, as well as higher twist functions. The transverse part of the metric tensor, denoted by $g^T_{\mu\nu}$ and the ϵ-tensor are contracted with the projection operators $\mathcal{P}_{\alpha\beta} = g_{\alpha\beta} - q_{2\alpha} q_{1\beta}/(q_1 q_2)$. Therefore, current conservation is manifest. Our definitions are chosen in such a way that in the forward case, i.e. $\Delta = 0$, the usual structure functions measured in DIS are

$$F_1(x_{\rm Bj}, Q^2) = \frac{1}{2\pi} {\rm Im} \mathcal{F}_1(\xi = x_{\rm Bj}, 0, Q^2, 0), \quad g_1(x_{\rm Bj}, Q^2) = \frac{1}{2\pi} {\rm Im} \mathcal{G}_1(\xi = x_{\rm Bj}, 0, Q^2, 0). \tag{5}$$

Since the virtuality of the incoming photon is deep in the Euclidean domain the hadron is probed almost on the light cone $x^2 \approx 0$. Hence, the amplitude can be straightforwardly treated by a non-local version of the light-cone operator product expansion [10]. More recently, it has been proven that assuming a smooth SPD the collinear singularities are indeed factorizable at leading twist to all orders of perturbation theory [6, 11, 12]. Therefore, the amplitudes $\mathcal{F}_1$ and $\mathcal{G}_1$ factorize in a perturbative hard scattering amplitude and a SPD $q(t, \eta, \Delta^2, \mu^2)$ as

$$\mathcal{T}^a = \sum_Q e_Q^2 \int_{-1}^{1} \frac{dt}{|\eta|} \left[{}^Q T^a(\xi, \eta, t, Q^2, \mu^2)\, {}^Q q^a(t, \eta, \Delta^2, \mu^2) + \frac{1}{N_f} \frac{1}{\eta} {}^G T^a(\xi, \eta, t, Q^2, \mu^2)\, {}^G q^a(t, \eta, \Delta^2, \mu^2) \right], \tag{6}$$

where $\mathcal{T}^V = \mathcal{F}_1$, $\mathcal{T}^A = \mathcal{G}_1$ and the sum runs over quark species $Q = u, d, s$ with electrical charge e_Q.

The non-perturbative input is concentrated in the SPDs which read in the Leipzig conventions [1]

$$\left\{ \begin{array}{c} {}^Q q^V \\ {}^Q q^A \end{array} \right\}(t, \eta) = \int \frac{d\kappa}{2\pi} e^{i\kappa t P_+} \langle P_2 S_2 | \bar{\psi}(-\kappa n) \left\{ \begin{array}{c} \gamma_+ \\ \gamma_+ \gamma_5 \end{array} \right\} \psi(\kappa n) | P_1 S_1 \rangle, \tag{7}$$

$$\left\{ \begin{array}{c} {}^G q^V \\ {}^G q^A \end{array} \right\}(t, \eta) = \frac{4}{P_+} \int \frac{d\kappa}{2\pi} e^{i\kappa t P_+} \langle P_2 S_2 | G^a_{+\mu}(-\kappa n) \left\{ \begin{array}{c} g_{\mu\nu} \\ i\epsilon_{\mu\nu-+} \end{array} \right\} G^a_{\nu+}(\kappa n) | P_1 S_1 \rangle. \tag{8}$$

For the quark distributions we have omitted for brevity the flavour indices. A form factor decomposition would give us the functions H, E and $\widetilde{H}$, $\widetilde{E}$ for the parity even and odd sectors, respectively, as introduced in Ref. [2].

Using the spatial parity and scaling properties evident from the convolution-type formulae of Ref. [13] we write the perturbative expansion of T as

$$\xi\ {}^{i}T^{a} = {}^{i}T^{a(0)}\left(\frac{\xi}{\eta}, \frac{t}{\eta}\right) + \frac{\alpha_s(\mu^2)}{2\pi}{}^{i}T^{a(1)}\left(\frac{\xi}{\eta}, \frac{t}{\eta}, \frac{Q^2}{\mu^2}\right) + \mathcal{O}\left(\alpha_s^2\right) \mp (t \to -t)\,, \tag{9}$$

where the "$-$"("$+$") sign corresponds to the $V(A)$ channel. The tree level coefficient functions are

$$ {}^{Q}T^{(0)}\left(\frac{\xi}{\eta}, \frac{t}{\eta}\right) = \frac{1}{1 - t/\xi - i\epsilon}, \qquad {}^{G}T^{(0)} = 0. \tag{10}$$

The hard scattering amplitude $T^{(1)}$ can be calculated by making use of standard methods of perturbative QCD [14]. At the same time there exists an interesting possibility to predict [13] these quantities with the help of the conformal operator product expansion (COPE). It has been introduced more than two decades ago in Ref. [15] and consequently applied to exclusive processes at leading order [16]. The main advantage of the COPE is that under the assumption of conformal covariance it predicts the Wilson coefficients of local conformal operators with a given conformal spin entering the expansion of the product of two currents up to a normalization constant. The latter is fixed by the known Wilson coefficients of forward DIS. Since the COPE is valid only for $|\xi| > 1$ we perform a summation over all conformal partial waves and obtain a representation of the hard scattering amplitudes as a convolution of kernels. E.g. for the non-singlet case we have

$$T(\xi, \eta, t, Q^2, \mu^2) = \frac{1}{\xi} F\left(\frac{\xi}{\eta}, \frac{r}{\eta}\right) \otimes \left(\frac{Q^2}{\mu^2}\right)^{V\left(\frac{r}{\eta}, \frac{s}{\eta}\right)} \otimes C\left(\frac{s}{\eta}, \frac{t}{\eta}; \alpha_s\right), \tag{11}$$

where the convolution is defined for any test functions $\tau_1(x, y)$ and $\tau_2(x, y)$ with $\tau_1 \otimes \tau_2 \equiv \int \frac{dy}{|\eta|} \tau_1(x, y)\tau_2(y, z)$. The evolution kernel $V\left(\frac{r}{\eta}, \frac{s}{\eta}\right)$ and the coefficient function $C\left(\frac{s}{\eta}, \frac{t}{\eta}; \alpha_s\right)$ are diagonal w.r.t. the conformal waves — the Gegenbauer polynomials:

$$\eta^j C_j^{3/2}\left(\frac{t}{\eta}\right) \otimes \begin{Bmatrix} V \\ C \end{Bmatrix} \left(\frac{t}{\eta}, \frac{t'}{\eta}; \alpha_s\right) = \begin{Bmatrix} -\frac{1}{2}\gamma_j(\alpha_s) \\ c_j(\alpha_s) \end{Bmatrix} \eta^j C_j^{3/2}\left(\frac{t'}{\eta}\right), \tag{12}$$

and their eigenvalues $\gamma_j(\alpha_s)$ and $c_j(\alpha_s)$ coincide with anomalous dimensions and Wilson coefficients, respectively, which appear in DIS. A closed form of the function F can be deduced order by order in α_s from the sum

$$F(x, y) = \sum_{j=0}^{\infty} \left(\frac{2x}{1+x}\right)^{j+1} \frac{B(j+1, j+2)}{(1+x)^{\gamma_j/2}} {}_2F_1\left(\left. \begin{matrix} 1 + j + \gamma_j/2,\ 2 + j + \gamma_j/2 \\ 4 + 2j + \gamma_j \end{matrix} \right| \frac{2x}{1+x}\right) C_j^{3/2}(y). \tag{13}$$

The expressions (11-13) which are exact in QCD up to NLO[2] define the so-called conformal subtraction (CS) scheme. To obtain the $\overline{\rm MS}$ results one has to perform a scheme transformation

[2]Beyond this order the trace anomaly appears and provides conformal noncovariant terms proportional to $\beta/g[\alpha_s/(2\pi) + O(\alpha_s^2)]$.

[13, 17], which is governed by a conformal anomaly appearing in special conformal Ward identities [18]. Finally, the NLO coefficient functions read

$$
{}^{Q}T^{(0)} = \frac{1}{1-t}, \qquad {}^{Q}T^{V(1)} = {}^{Q}T^{A(1)} - \frac{C_F}{1+t} \ln \frac{1-t}{2}, \tag{14}
$$

$$
{}^{Q}T^{A(1)} = \frac{C_F}{2(1-t)} \left[\left(2 \ln \frac{1-t}{2} + 3 \right) \left(\ln \frac{-q_1^2}{\mu^2} + \frac{1}{2} \ln \frac{1-t}{2} - \frac{3}{4} \right) - \frac{27}{4} - \frac{1-t}{1+t} \ln \frac{1-t}{2} \right], \tag{15}
$$

$$
{}^{G}T^{V(1)} = -{}^{G}T^{A(1)} + \frac{N_f}{2} \left[\frac{1}{1-t} \left(\ln \frac{-q_1^2}{\mu^2} + \ln \frac{1-t}{2} - 2 \right) + \frac{\ln \frac{1-t}{2}}{1+t} \right], \tag{16}
$$

$$
{}^{G}T^{A(1)} = \frac{N_f}{2} \left[\left(\frac{1}{1-t^2} + \frac{\ln \frac{1-t}{2}}{(1+t)^2} \right) \left(\ln \frac{-q_1^2}{\mu^2} + \ln \frac{1-t}{2} - 2 \right) - \frac{\ln^2 \frac{1-t}{2}}{2(1+t)^2} \right]. \tag{17}
$$

We have used here the scaling property mentioned above and the fact that the hard scattering amplitude for DVCS, i.e. $\xi = \eta$, is a simple analytical continuation of the ones for the production of a scalar and pseudo-scalar mesons, respectively, in the collision of a real and a highly virtual photon.

To the same accuracy we have to include the two-loop evolution of the SPDs. Based on the knowledge of NLO anomalous dimensions constructed from conformal constraints [18] we will use the methods of Ref. [9] where we have already dealt with scaling violations in NLO.

For numerical estimates of the DVCS amplitude we expand the hard scattering amplitude in terms of Legendre polynomials and reexpress the expansion coefficients in terms of conformal moments. This allows us to include easily the NLO evolution of the SPDs as described in Ref. [9]. Unfortunately, this method is reliable only for moderate values of $x_{\rm Bj}$ since the convergence of the series of polynomials is not under control for very low $x_{\rm Bj}$. Already for $x_{\rm Bj} \sim 0.05$ one needs about 180 polynomials.

To model the SPD one can first explore the simplest possibility of equating it to the usual forward parton density. We will designate this as FPD-model. For small η, where the so-called DGLAP region with $|t| > |\eta|$ dominates, this may be justified. However, it is an open problem how the exclusive region $|t| < |\eta|$ corresponding to the production or absorption of a meson-like state looks like. To get a more realistic model one can use the relation of SPDs to the so-called double distributions (DDs) $f(z_-, z_+)$, introduced in [1] and rediscovered in [6]

$$
q(t, \eta, Q^2) = \int_{-1}^{1} dz_+ \int_{-1+|z_+|}^{1-|z_+|} dz_- \delta(z_+ + \eta z_- - t) f(z_-, z_+, Q^2). \tag{18}
$$

The latter, according to Ref. [19], is given by the product of a forward distribution $f(z)$ (more precisely $q(z)$ for quarks and $zg(z)$ for gluons) with a profile function π

$$
f(z_-, z_+) = \pi(z_-, z_+) f(z_+), \tag{19}
$$

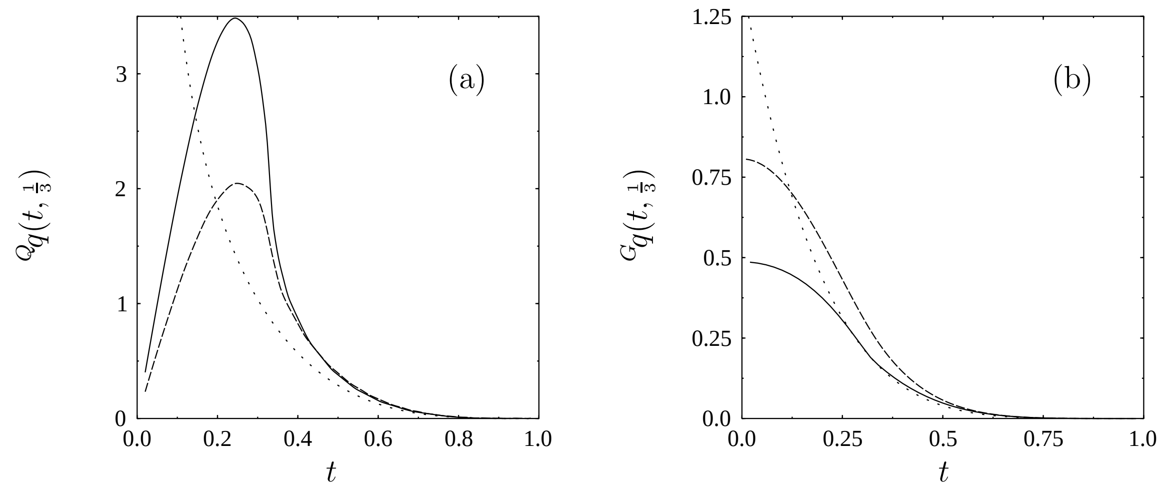

Figure 1: In (a) and (b) we show the SPD for the sum of the non-polarized u and $\bar{u}$ quarks and the gluon densities, respectively. The FPD-model (dotted line) and the DD-model with MRS [20] parametrization (dashed line) are modeled at a scale of 4 GeV^2, while the DD-Model with GRV [21] parametrization (solid line) has been taken at the momentum scale 0.4 GeV^2 and has been evolved afterwards with NLO formulae to $\mathcal{Q}_0^2 = 4$ GeV^2. The skewedness parameter is $\eta = 1/3$ which corresponds to $x_{\rm Bj} = 1/2$.

where π for quarks and gluons is given by

$$ {}^Q\pi(z_-, z_+) = \frac{3}{4}\frac{[1-|z_+|]^2 - z_-^2}{[1-|z_+|]^3}, \qquad {}^G\pi(z_-, z_+) = \frac{15}{16}\frac{\left\{[1-|z_+|]^2 - z_-^2\right\}^2}{[1-|z_+|]^5}, \tag{20} $$

respectively. We will refer to this prescription as the DD-model. Note, that parton distributions used here have the support $-1 \le z \le 1$ and, therefore, the contributions $-q(-z)$, $+\Delta q(-z)$ with $z \ge 0$ have the usual interpretation as antiquarks.

Now we are in a position to study the NLO amplitudes in the $\overline{\rm MS}$ scheme for the models specified above. In what follows we use the convention $\mathcal{Q}^2 \equiv -q_1^2$ and equate the factorization scale with $-q_1^2$. The value of $\Lambda_{\overline{\rm MS}}$ is $\Lambda^{(4)}_{\rm NLO} = 246$ MeV. Since we will rely on available parametrizations of the parton densities which are defined at different scales the DD-models will differ as well. Given a forward distribution at an input scale e.g. $\mathcal{Q}_0^2 = 4$ GeV^2, as e.g. for the MRS fit [20], it can be viewed as evolved only according to the DGLAP equation from a lower scale and then folded with a profile to form the DD-model. If the former is defined at an input of $\mathcal{Q}_0^2 = 0.4$ GeV^2, like e.g. the GRV densities [21], to confront both models we have to evolve GRV-based SPDs with non-forward evolution equations. We demonstrate these features of the DD-models in Fig. 1 where we have compared the MRSS0 [20] with the GRV parametrization [21] according to the procedure sketched above.

In general we have found that the amplitudes calculated from the models mentioned above are similar in shape. However, in the parity even sector the DVCS amplitude is very sensitive to

the $z \to 0$ behaviour of the sea quarks. Thus, different models will produce qualitatively different predictions. This is demonstrated in Fig. 2(a) for $\mathcal{F}_1$, where even the sign of the amplitude for small $x_{\rm Bj}$ differs for the FPD- and DD-model. Note that the sign of the FPD will change for very low $x_{\rm Bj}$ and that a slightly different parametrization gives a prediction similar to the one for a DD-model. We see from this figure that the NLO corrections are as large as 50% and even more. Evidently, the imaginary part is more sensitive to the shape of the model distribution than the real part. Their ratio is shown in Fig. 2(b). In Fig. 2(c) we compare the predictions of the DD-model for the MRS and the GRV parametrizations taken at different input scales as explained above. This affects the size of both the real and imaginary part, so that the ratio only slightly changes with the scale as shown in Fig. 2(d). The ratio is also not sensitive to the radiative corrections in the coefficient functions. Note that the evolution will suppress the DGLAP region, so that the asymptotic distributions are concentrated in the region $|t| \le \eta$ and are given by the terms with the lowest conformal spin in the conformal expansion of SPD. Thus for asymptotically large $\mathcal{Q}^2$ the imaginary part has to vanish and the radiative corrections to the real part are determined by the lowest Wilson coefficient in the COPE.

In Fig. 3 we show the predictions for $\mathcal{G}_1$ with the GSA parametrization [22] and investigate the size of radiative corrections in detail. In this case we have found a similar model dependence as in the previous case: predictions are sensitive to the sea quark parametrization which turn on at $x_{\rm Bj} \sim 0.1$ and also may cause very substantial radiative corrections. Since the polarized sea quark distributions are not well known, we simply neglect them. The real and imaginary parts for the DD-model are shown in Fig. 3(a,b), respectively, at the input scale 4 GeV^2 and evolved upwards to the scale 10 GeV^2. Not surprisingly the size of the radiative corrections is decreasing with increasing $\mathcal{Q}^2$. Note that for large $x_{\rm Bj}$ the ratio of real to imaginary part provides again a useful information about the SPD as discussed above. The size of radiative corrections for $\mathcal{F}_1$ and $\mathcal{G}_1$ are similar. In Fig. 3(c,d) we demonstrate the factorization scale dependence, where we took in the hard scattering amplitude and SPD the scale to be $\mu^2 = \{\mathcal{Q}^2/2, \mathcal{Q}^2, 2\mathcal{Q}^2\}$. We have also studied the perturbative corrections in the conformal subtraction scheme, however, we have found that in the DVCS kinematics this scheme only slightly reduces the NLO corrections.

To summarize, we have studied the DVCS amplitude in perturbative QCD for the region $x_{\rm Bj} \ge 0.05$. It turns out that a separate measurement of real and imaginary part in this kinematic domain would provide us with information about the SPDs, which is indispensable to discriminate between different models. Moreover, the amplitude at $x_{\rm Bj} \lesssim 0.1$ depends in a crucial way on both the model and the used parametrization of the forward parton distributions, especially, of the sea distribution. Generally, we found that the radiative correction can be as large as 50% and more, depending on $x_{\rm Bj}$. Fortunately, the ratio of real to imaginary part is less sensitive to these

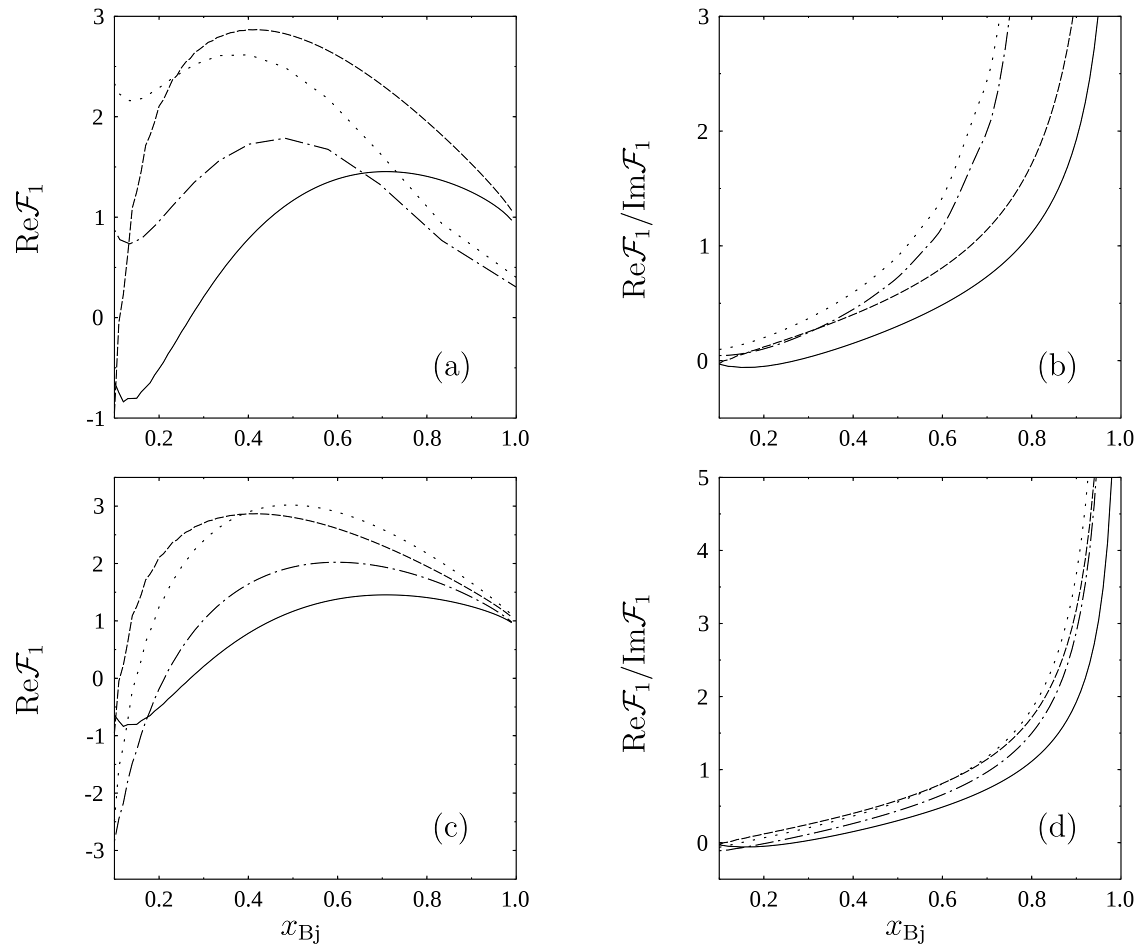

Figure 2: The real part (a) and the ratio of the real to imaginary part (b) of $\mathcal{F}_1$, respectively, for the DD- and FPD-model with the MRS parametrization taken at the input scale $\mathcal{Q}_0^2 = 4$ GeV^2. Here the DD(FPD)-model at LO and NLO is shown as dashed (dotted) and solid (dash-dotted) lines, respectively. Plots (c) and (d) show the same as in (a) and (b), respectively but for MRS and GRV parametrizations. Here MRS (GRV) amplitude is plotted at LO as dashed (dotted) and at NLO as solid (dash-dotted) lines, respectively.

radiative corrections and is only mildly scheme dependent. Thanks to this property it can give us more insights into the structure of SPDs.

The kinematical situation relevant for HERA experiments requires an extension of our analysis downwards to very low x_{Bj}. In this case the polynomial method used here will not be the most appropriate tool. Fortunately, a direct numerical convolution of the hard scattering part and skewed parton distributions can be carried out without major difficulties. The remaining problem is to evolve the SPD in this kinematics, which may be solved by numerical integration of the evolution equation. Recently, the evolution kernels at two-loop order have been completely constructed [23] from the knowledge of conformal anomalies and splitting functions and extensive use of supersymmetric constraints. This provides in future the opportunity to study evolution

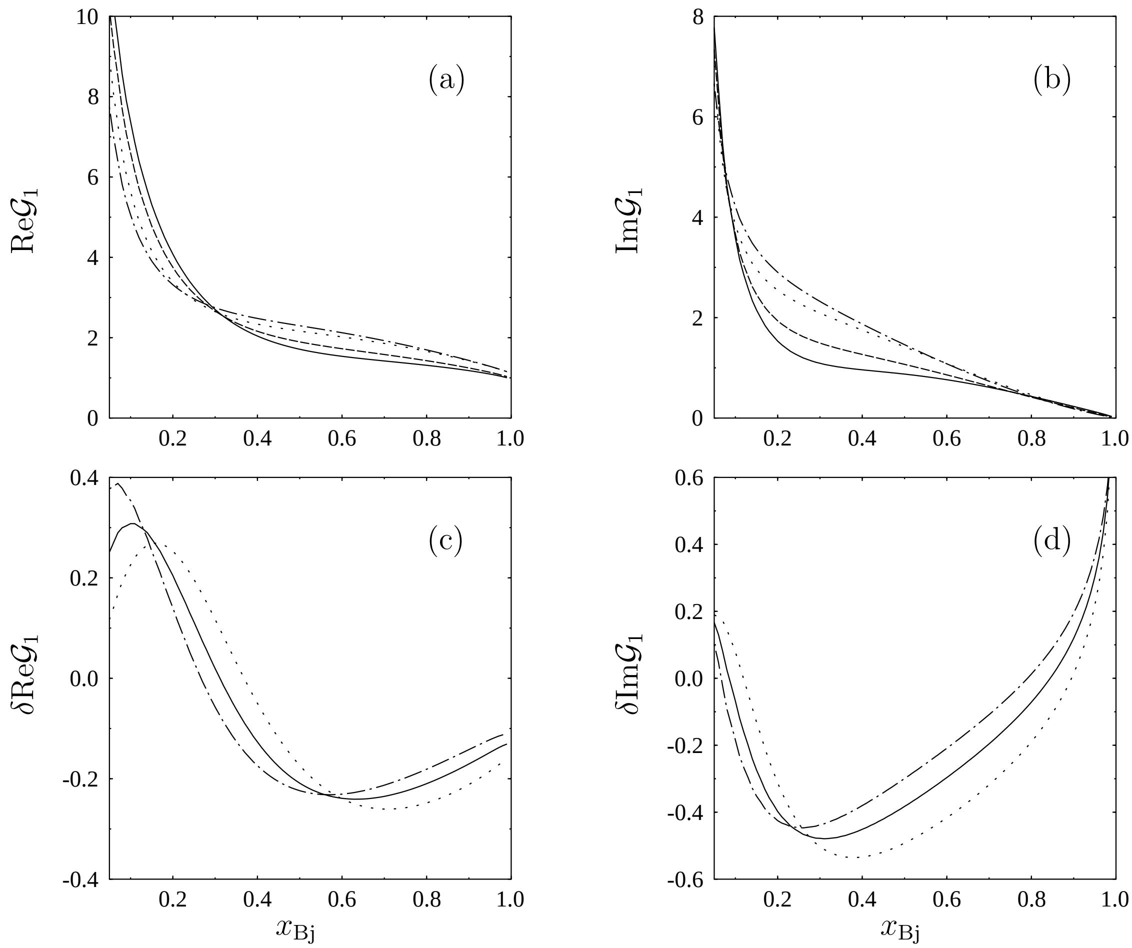

Figure 3: The real (a) and the imaginary (b) part of $\mathcal{G}_1$ for the DD-model with the GS [22] parametrization taken at the input scale $\mathcal{Q}_0^2 = 4$ GeV^2 [LO (dotted line) and NLO (solid line)] as well as evolved to $\mathcal{Q}^2 = 10$ GeV^2 [LO (dash-dotted line) and NLO (dashed line)]. Plot (c) and (d) shows the relative radiative corrections $\delta = [(\mathrm{NLO} - \mathrm{LO})/\mathrm{LO}]$ again for the real and imaginary part, respectively, for $\mathcal{Q}^2 = 4$ GeV^2. The factorization scale μ^2 is set equal to $\mathcal{Q}^2/2$ (dotted line), $\mathcal{Q}^2$ (solid line), and $2\mathcal{Q}^2$ (dash-dotted line).

effects also at very low $x_{\rm Bj}$ which is a very important task in view of ongoing experiments on diffractive meson production at HERA.

To decrease the theoretical uncertainties due to radiative corrections, it is necessary to go beyond NLO. In a first step one can include only the hard-scattering amplitude to two-loop order accuracy. Of course, a direct calculation will be very cumbersome. Fortunately, however, in the conformal limit of the theory, when the QCD Gell-Mann–Low function formally equals zero, we can rely directly on the COPE and it seems feasible to extend this technology to NNLO. Conformal symmetry breaking effects can be perturbatively calculated and one piece of information is already available [24]. From the practical point of view it should be stressed that all high hopes connected to SPDs rely on the assumption that the higher order terms are controllable. We want to emphasize

that our approach allows to study efficiently the structure of perturbative corrections and might be crucial to establish the same rigour of treatment as for the forward parton densities.

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